Counter-polarized single-photon generation from the auxiliary cavity of a weakly nonlinear photonic molecule

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We propose a scheme for the resonant generation of counter-polarized single photons in double asymmetric cavities with a small Kerr optical nonlinearity (as that created by a semiconductor quantum well) compared to the mode broadening. Due to the interplay between spatial intercavity tunneling and polarization coupling, by weakly exciting with circularly polarized light one of the cavities, we predict strong antibunching of counter-polarized light emission from the non-pumped auxiliary cavity. This scheme due to quantum interference is robust against surface scattering of pumping light, which can be suppressed both by spatial and polarization filters.

Single photons play an essential role in quantum information technologies and their generation is a fascinating subject in the fields of quantum optics and condensed matter physics. As deterministic (on-demand) sources of single photons, single semiconductor quantum dots embedded in optical microcavities have been investigated with impressive results in the case of non-resonant excitation.[1–6] However, if one wishes to build an array of single-photon emitters with intentional pattern on the same wafer, the completely random position of self-organized quantum dots is by definition a kind of limitation. In the case of non-resonant generation, the repetition rate of the single-photon source is limited by the relatively slow relaxation time of the injected carriers in the semiconductor device. For higher repetition rates and freedom in array design, the resonant photon blockade in photonic pillars including a Kerr medium has been proposed, where the nonlinear medium is a simple semiconductor quantum well,[7] which can be patterned with great flexibility. Unfortunately, the non-trivial difficulty of such approach is that the strength of the Kerr nonlinearity must be much larger than the mode broadening, thus requiring in practice the use of ultrasmall photonic resonators with very high quality factors.

In a recent letter,[8] it has been proved that by using two coupled pillars (a double cavity system or photonic 'molecule') it is possible to get very pure single-photon emission with a nonlinearity surprisingly small with respect to the losses. As analytically shown in Ref. 9, the antibunching originates from the destructive quantum interference between excitation and tunneling paths of photons. However, a practical hurdle still remains: in the proposed scheme,[8,9] single photons are emitted from the excited cavity with the same polarization of the pump. Hence, spurious effects as pump surface scattering could mask such an effect. In the present paper, we propose another scheme based on the generalization of the photonic molecule approach by taking advantage of the polarization degree of freedom in asymmetric cavities having a frequency splitting between modes with orthogonal linear polarizations. We show how to get counter-polarized single-photons from the non-pumped auxiliary cavity, thus providing a way to get rid of the pump spurious scattering by both spatial and polarization filtering.

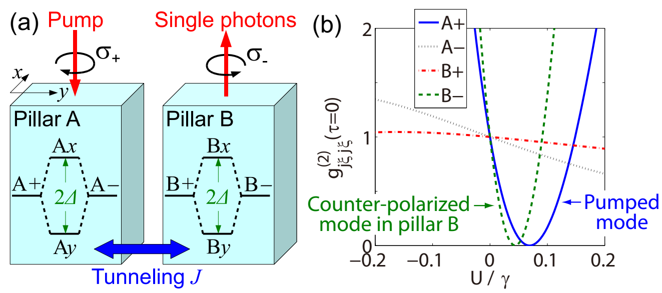

FIG. 1. (a) Sketch of the system consisting of two coupled micropillars: due to the shape anisotropy, the photon eigenmodes have orthogonal linear polarizations. The mode with ξ polarization in pillar j is denoted as $j\xi$, and the energy levels for the single photon states in each pillar are depicted. By illuminating circularly polarized light on pillar A, counter-polarized single photons are emitted from pillar B even with a small nonlinearity. (b) The equal-time second-order correlation functions $\{g^{(2)}_{j\xi j\xi}(\tau=0)\}$ are plotted as functions of nonlinearity U normalized to broadening γ. The tunneling strength is $J = 5\gamma$, the polarization splitting is $\Delta = 2.5\gamma$, and the pump frequency is tuned as $\delta E = E - \hbar\omega_{\mathrm{p}} = 0.2772\gamma$. In addition to the antibunching of the pumped mode $A+$, nearly perfect antibunching is obtained in mode $B-$ with the relatively small nonlinearity $U = 0.0438\gamma$.

As a realistic system, we consider two spatially separated semiconductor micropillars with asymmetric shape,[10] coupled with a photonic tunneling strength J. Each pillar has two different linearly polarized photonic modes (x and y directions), energetically split by 2Δ due to the shape anisotropy. Fig. 1(a) shows a sketch of con-

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sidered system. The Hamiltonian is represented as

$$\hat{H} = \sum_{j=\{A,B\}} \left[(E+\Delta)\hat{a}^{\dagger}_{jx}\hat{a}_{jx} + (E-\Delta)\hat{a}^{\dagger}_{jy}\hat{a}_{jy} \right] + \sum_{\xi=\{x,y\}} J(\hat{a}^{\dagger}_{A\xi}\hat{a}_{B\xi} + \mathrm{H.c.}) + (F\mathrm{e}^{-\mathrm{i}\omega_{\mathrm{p}}t}\hat{a}^{\dagger}_{A+} + \mathrm{H.c.}) + \sum_{j=\{A,B\},\xi=\{+,-\}} U\hat{a}^{\dagger}_{j\xi}\hat{a}^{\dagger}_{j\xi}\hat{a}_{j\xi}\hat{a}_{j\xi}. \quad (1)$$

Here, $\hat{a}_{j\xi}$ is the annihilation operator of a photon with polarization ξ in pillar j. The relation between circularly and linearly polarized modes is given by the standard operator expression $\hat{a}_{j\pm} = (\hat{a}_{jx} \pm \mathrm{i}\hat{a}_{jy})/\sqrt{2}$. We consider the configuration where pillar A is pumped with σ_+-circular polarization, being ω_{p} and F the pump frequency and amplitude respectively. The pumping strength is moderate to guarantee the average number of photons in the system not exceeding unity. If the average number is increased, antibunching is worsened, because the quantum interference in the present scheme is valid if three-photon subspace can be neglected, as in the previously considered scenario.[8,9] The nonlinearity is represented by the last term in Eq. (1) conserving the total spin of two photons. The effective nonlinearity can be mediated by the presence of a quantum well excitonic resonance,[11] but this effective Kerr term is quite general for systems with a third-order nonlinearity. The cross-polarized term such as $U_{\mathrm{cross}}\hat{a}^{\dagger}_{j+}\hat{a}^{\dagger}_{j-}\hat{a}_{j-}\hat{a}_{j+}$ is not considered in the present paper, because U_{cross} is usually much smaller than U,[12] but it could be added without qualitative changes (not shown). By using the theoretical method detailed in Ref. 7, we have numerically calculated second-order correlation functions $g^{(2)}_{j\xi j'\xi'}(\tau) = \langle\hat{a}^{\dagger}_{j\xi}\hat{a}^{\dagger}_{j'\xi'}(\tau)\hat{a}_{j'\xi'}(\tau)\hat{a}_{j\xi}\rangle/\langle\hat{a}^{\dagger}_{j\xi}\hat{a}_{j\xi}\rangle\langle\hat{a}^{\dagger}_{j'\xi'}\hat{a}_{j'\xi'}\rangle$ at the steady state under continuous pumping and a dissipation of photons with a rate $\gamma/\hbar$ in each mode.

In Fig. 1(b), we plot equal-time correlations $\{g^{(2)}_{j\xi j\xi}(\tau = 0)\}$ as a function of nonlinearity U normalized to γ. We consider the tunneling strength $J = 5\gamma$, polarization splitting $\Delta = 2.5\gamma$, and pump detuning $\delta E = E - \hbar\omega_{\mathrm{p}} = 0.2772\gamma$. In addition to antibunching at the pumped mode $A+$ due to the previously proposed scheme[8,9] strong antibunching is achieved for mode $B-$ with the (small) optimal nonlinearity $U = 0.0438\gamma$. In the weak pumping limit, $g^{(2)}_{B-B-}(\tau = 0)$ is reduced to zero. The underlying destructive quantum interference is different from the previous one in Refs. 8 and 9. By deriving the equations of motions for the amplitudes of the possible Fock states for the zero-, one- and two-photon states (generalizing the method in Ref. 9) we have derived the optimal conditions of the counter-polarized antibunching and found the interference paths leading to the antibunching in mode $B-$. In Fig. 2(a), the zero-, one-, and two-photon state manifolds are depicted and labeled as $|0\rangle$, $|j\xi\rangle$, and $|j\xi, j'\xi'\rangle$, respectively. The paths responsible to antibunching in mode $|B-, B-\rangle$ are the J-assisted (spatial tunneling) path from $|A-, B-\rangle$ and the

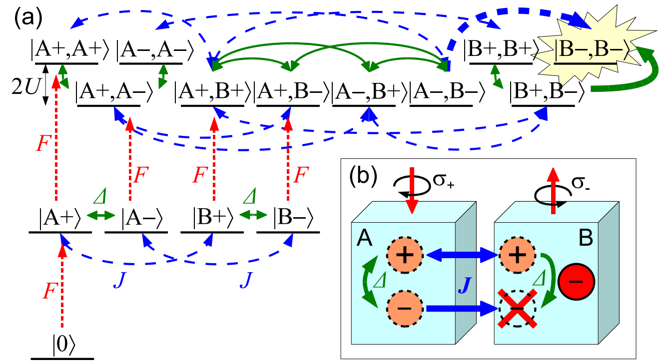

FIG. 2. (a) Sketch of all the transition paths between zero-photon $|0\rangle$, one-photon $|j\xi\rangle$, and two-photon states $|j\xi, j'\xi'\rangle$. The antibunching of mode $B-$ is due to the destructive quantum interference between the two paths from $|A-, B-\rangle$ and $|B+, B-\rangle$ to $|B-, B-\rangle$. (b) Pictorial representation. If there is already a "−" photon in pillar B, another "−" photon cannot exists in the same pillar because of the interference between the J-assisted (spatial tunneling) path from "−" photon in pillar A and the Δ-assisted (polarization coupling) path from "+" photon in pillar B. This quantum interference occurs for an optimal value of the nonlinearity and laser detuning.

Δ-assisted (polarization coupling) one from $|B+, B-\rangle$: destructive interference occurs for an optimal value of the nonlinearity U, which is much smaller than the inverse of the cavity lifetime. Fig. 2(b) shows a pictorial interpretation. In presence of "−"-photon in cavity B, another photon cannot enter the same pillar due to the quantum interference. The presence of a cross-polarized nonlinear term U_{cross} would not change this picture (not shown) and it does not create a new path to $|B-, B-\rangle$. Furthermore, while identical pillars are supposed in the reported calculation, we have numerically checked that nearly perfect antibunching in mode $B-$ can be obtained even under a deviation of the order of γ on eigen frequencies E and splitting Δ between the two pillars, and also tunneling strength J between two polarizations, while the pumping frequency ω_{p} should be properly tuned.

From the equations of motions of up to the two-photon Fock states, we have calculated the optimal nonlinearity U_{opt} for the perfect antibunching under the weak pumping limit. Fig. 3(a) shows U_{opt} as a function of Δ/γ, and Fig. 3(b) represents the ratio between the average photon number $n_{B-} = \langle\hat{a}^{\dagger}_{B-}\hat{a}_{B-}\rangle$ in mode $B-$ and the total number $n_{\mathrm{total}} = \sum_{j,\xi}\langle\hat{a}^{\dagger}_{j\xi}\hat{a}_{j\xi}\rangle$ for the corresponding U_{opt}. The minimum nonlinearity that is required for the antibunching is decreased by the increase of J, and the required nonlinearity is decreased together with the splitting Δ (or the opposite behavior when $J < \Delta$). However, together with the decrease of U_{opt}, the occupation probability of mode $B-$ is significantly decreased as seen in Fig. 3(b). For $U_{\mathrm{opt}} < 0.1\gamma$, the curve in Fig. 3(b) is almost saturated at $J = 5\gamma$, and we can obtain a probability $n_{B-}/n_{\mathrm{total}} \sim 10^{-2}$, which corresponds to a gener-

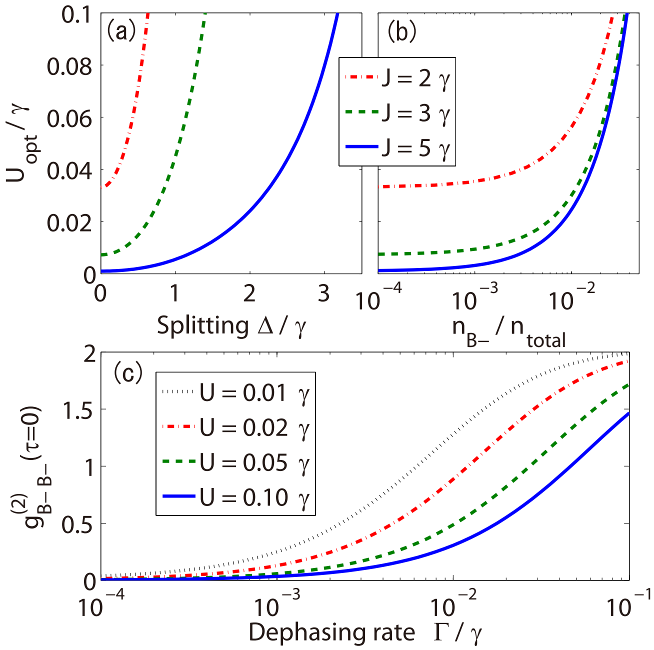

FIG. 3. (a) Optimal nonlinearities $U_{\rm opt}$ are plotted as a function of Δ/γ for tunneling strengths $J/\gamma = 2$, 3, and 5. (b) Under the optimal conditions, the ratios between the average number n_{B-} of photons in mode $B-$ and the number $n_{\rm total}$ in the total system are plotted for corresponding $U_{\rm opt}/\gamma$. (c) The obtainable $g^{(2)}_{B-B-}(\tau = 0)$ are plotted versus the dephasing rate $\mathit{\Gamma}$ normalized to γ. The tunneling strength is $J = 5\gamma$, energy splitting $\mathit{\Delta}$ and pumping frequency are chosen to give the nearly perfect antibunching for each nonlinearity U in the absence of pure dephasing.

ation rate of the order of 100 MHz for a cavity lifetime in the picosecond rage. This rate is higher than that of the quantum dots[1–6] by one order of magnitude.

Finally, we have examined the robustness of the present scheme against dephasing of photons. Since quantum interferences are responsible in the present and previous schemes,[8,9] pure dephasing can decrease the quality of the antibunching. By using the standard pure dephasing model due to quadratic coupling with a reservoir,[13] we consider dephasing with a rate $\mathit{\Gamma}/\hbar$ affecting linearly polarized modes of each pillar (the results shown below are not significantly modified even if the dephasing is supposed to affect the circularly polarized modes). Fig. 3(c) shows $g^{(2)}_{B-B-}(\tau = 0)$ as a function of $\mathit{\Gamma}/\gamma$. The tunneling strength is $J = 5\gamma$, and the splitting $\mathit{\Delta}$ is chosen to give perfect antibunching for each nonlinearity U in the absence of the dephasing. As clearly shown, the antibunching can be significantly worsened in presence of pure dephasing for a given value of the nonlinearity. However, even if the nonlinearity is quite small, for example $U = 0.05\gamma$, antibunching is still observable if $\mathit{\Gamma} = 10^{-2}\gamma$ and becomes very strong if $\mathit{\Gamma} = 10^{-3}\gamma$. The curves in Fig. 3(c) does not strongly depend on the tunneling strength J if the nonlinearity U is large enough compared to the corresponding minimum shown in Fig. 3(a). Since the optimal nonlinearity can be weak in the present scheme, one can consider cavities with relatively small photon lifetime (small quality factor) in a regime where the pure dephasing time can be thus neglected, a very promising outlook.

In conclusion, we have proposed a scheme of single-photon generation due to a destructive quantum interference effect in a weakly nonlinear double cavity system, where each cavity have two linearly polarized, frequency-split modes. Due to the spatial tunneling between the two cavities and the coupling between opposite circular polarizations, we have found that strong antibunching of counter-polarized emission can be obtained at the non-pumped auxiliary pillar. This new effect is of practical implications, because it provides a direct way to suppress the pump scattering via spatial and polarization filters. This scheme can also exported to arrays of nonlinear photonic molecules.

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[1] P. Michler, A. Kiraz, C. Becher, W. V. Schoenfeld, P. M. Petroff, L. Zhang, E. Hu, and A. Imamoglu, Science **290**, 2282 (2000).
[2] E. Moreau, I. Robert, L. Manin, V. Thierry-Mieg, J. M. Gérard, and I. Abram, Phys. Rev. Lett. **87**, 183601 (2001).
[3] J. Vuckovic, D. Fattal, C. Santori, G. S. Solomon, and Y. Yamamoto, Appl. Phys. Lett. **82**, 3596 (2003).
[4] D. Press, S. Götzinger, S. Reitzenstein, C. Hofmann, A. Löffler, M. Kamp, A. Forchel, and Y. Yamamoto, Phys. Rev. Lett. **98**, 117402 (2007).
[5] A. J. Shields, Nat. Photon. **1**, 215 (2007).
[6] J. Claudon, J. Bleuse, N. S. Malik, M. Bazin, P. Jaffrennou, N. Gregersen, C. Sauvan, P. Lalanne, and J.-M. Gerard, Nat. Photon. **4**, 174 (2010).
[7] A. Verger, C. Ciuti, and I. Carusotto, Phys. Rev. B **73**, 193306 (2006).
[8] T. C. H. Liew and V. Savona, Phys. Rev. Lett. **104**, 183601 (2010).
[9] M. Bamba, A. Imamoğlu, I. Carusotto, and C. Ciuti, Phys. Rev. A **83**, 021802(R) (2011).
[10] D. Bajoni, P. Senellart, E. Wertz, I. Sagnes, A. Miard, A. Lemaître, and J. Bloch, Phys. Rev. Lett. **100**, 047401 (2008).
[11] C. Ciuti, V. Savona, C. Piermarocchi, A. Quattropani, and P. Schwendimann, Phys. Rev. B **58**, 7926 (1998).
[12] A. Amo, S. Pigeon, C. Adrados, R. Houdré, E. Giacobino, C. Ciuti, and A. Bramati, Phys. Rev. B **82**, 081301 (2010).
[13] D. F. Walls, M. J. Collet, and G. J. Milburn, Phys. Rev. D **32**, 3208 (1985).